

Quantum creep in layered antiferromagnetic superconductor

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Abstract. In the mixed state of layered superconductor the antiferromagnetic order of magnetic ions could possibly create the spin-flop domains along the phase cores of the Josephson vortices. We discuss how this feature affects the activation energy in the macroscopic quantum tunnelling of the Josephson vortex. We show that the action and hence activation energy is rendered temperature dependent so that the quantum tunnelling rate becomes temperature dependent below the crossover temperature. We show that in such systems in constant temperature it is possible to occur of two regimes of the flux creep, thermal or quantum, dependent on the direction or intensity of the applied magnetic field.

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1. Introduction

The discoveries of ternary Rare Earth (RE) Chevrel Phases REMo_6S_8 and RERh_4B_4 [1] compounds with regular distribution of localized magnetic moments of RE atoms have proved conclusively the coexistence of various types of magnetism with superconductivity. Intensive experimental and theoretical research has shown that 4f electrons of RE atoms responsible for magnetism and 4d electrons of molybdenum chalcogenide or rhodium boride clusters responsible for superconductivity are spatially separated and therefore their interaction is weak. In many of these systems superconductivity coexists rather easily with antiferromagnetic order, where usually the Neel temperature T_N is lower than the critical temperature for superconductivity T_c . Quite often the ferromagnetic order is transformed into a spiral or domain-like structure, depending on the type and strength of magnetic anisotropy in the system [2, 3]. For almost two decades the problem of the interaction between magnetism and superconductivity has been overshadowed by high temperature superconductivity (HTS) found in copper oxides. However, the discovery of the magnetic order in Ru-based superconductors [4, 5, 6, 7] inspired a return to the so-called coexistence phenomenon [7]. The interplay between magnetism and superconductivity was studied in d-electron UGe_2 [8] and ZrZn_2 [9], where itinerant ferromagnetism may coexist with

superconductivity, and in heavy fermion UPd₂Al₃[14], where magnetic excitons are present in the superconducting phase. The recent discovery of the iron pnictide superconductors [15] have triggered broad interest in the mechanism of the coexistence of magnetism and superconductivity in this new class of superconductors [16]. They exhibit qualitative similarity to cuprates in that superconductivity occurs upon carrier doping (electrons, in this case) of pristine compounds that exhibit magnetism [17].

Among classic magnetic superconductors, the Chevrel phases have been studied most intensively. These compounds are mainly polycrystalline materials. However, some specific features can be measured only on single crystals. One such effect is a two-step flux penetration process, predicted in Ref.[18, 19] and later observed solely in the antiferromagnetic superconductor (bct) ErRh₄B₄ [20]. This very interesting phenomenon was recently rediscovered in DyMo₆S₈ [23], even though good quality single crystals of the classic antiferromagnetic superconductors have been measured for years. The specific feature caused by the long-range antiferromagnetic order in the mixed state of superconductor might be the creation of spin-flop (or metamagnetic) domain along each vortex core.

Consider an antiferromagnet with two magnetic sublattices as an example. An infinitesimal magnetic field applied perpendicular to the easy axis makes the ground antiferromagnetic state unstable against the phase transformation to the canted phase (spin-flop). On the contrary, if the magnetic field is applied parallel to the easy axis the antiferromagnetic configuration is stable up to the thermodynamic critical field H_T . When the field is further increased a canted phase develops in the system. Assume that in the antiferromagnetic superconductor the lower critical field fulfils the relation $H_{c1} < \frac{1}{2}H_T$ and that the external field, $H_{c1} < H < \frac{1}{2}H_T$, is applied parallel to the easy axis. Then the superconducting vortices appear in the ground antiferromagnetic state. If the field is increased above $\frac{1}{2}H_T$ the phase transition to the canted phase originates in the vortex core because the field intensity in the core doubles the external one. The spatial distribution of the field around the vortex is a decreasing function of the distance from its center. Hence the magnetic field intensity in the neighborhood of the core is less than H_T . Therefore, the rest of the vortex remains in the antiferromagnetic configuration. The radius of spin flop domain grows as the field is increased.

Thus, in the considered model there are two distinct types of vortices. Possible candidate of such system might be *ErBa₂Cu₃O₇* [10, 11]. This compound has tetragonal unit cell with small orthorhombic distortion in the *ab* plane. The *Er* ions form two sublattices antiferromagnetic structure of magnetic moments laying parallel and antiparallel to the **a** direction. The theoretical structure used in the present calculations consists of superconducting layers of thickness d_s and isolating ones of thickness d_i , $d = d_i + d_s$. In the isolating layers, the magnetic moments are running parallel and antiparallel to the **a** direction (easy axis). The magnetic field aligned parallel to the conducting planes makes the vortex lattice accommodate itself to the layer structure so that the vortex cores lie in between the superconducting sheets.

The structure of a vortex lying in the *ab* plane in a layered superconductor with

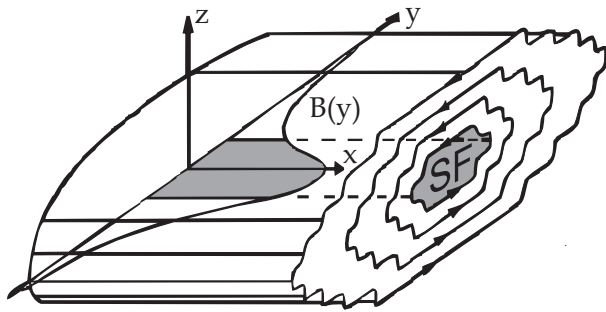


Figure 1. Single Josephson vortex lying in the ab plane along the \hat{x} -axis (\hat{a} -axis). The SF domain induced along the phase core is shown in the gray area.

Josephson coupling between adjacent layers resembles the Abrikosov's one except that the order parameter does not vanish anywhere [30]. Instead there exists a region, r_j along the plane and d perpendicular to it, where the Josephson current j_z is of the order of the critical current. In this region, which plays the role of the vortex core, the London model fails. Away from the core the streamlines of the shielding supercurrents, which also represents contours of constant magnetic field, are elliptical except for the zigzags, shown in Fig.(1), due to the intervening insulating layers.

2. Macroscopic quantum tunnelling through intrinsic pinning barriers

A current density j , flowing along the planes perpendicular to the applied magnetic field exerts a Lorentz force on the vortices in the \mathbf{c} direction so that intrinsic pinning barriers are formed on strongly superconducting layers. The experimental evidence of quantum tunnelling is based on the fact that the magnetic moment relaxation rate exhibits two types of behavior as a function of temperature. Above a characteristic temperature T_0 in the thermal activation regime the decay rate is of the Arrhenius type $\Gamma \sim \exp(-U_0/k_B T)$. Below T_0 , the decay rate is assumed *a priori* to be essentially independent of temperature $\Gamma \sim \exp(-S/\hbar)$ and is interpreted as arising from the quantum tunneling of vortices through intrinsic pinning potential [25, 26, 27]. In the following we show a considerable change of tunneling rate and crossover temperature due to the SF phase formation in the vortex core.

2.1. General formulae

We shall assume that the vortex line is a straight string-like object of effective mass m per unit length moving in a metastable intrinsic pinning potential $V(u)$ and exposed to continuous deformation $u(x, t)$ in the \hat{z} direction. The magnetic field is applied in \hat{x} (\hat{a} -axis) direction. Following Caldeira and Leggett [12] the vortex is coupled to a heat-bath reservoir assumed to be representable as a set of harmonic oscillators interacting linearly with vortex. In the semi-classical approximation the quantum decay rate is calculated

as a saddle-point solution (bounce) of the Euclidean action S for the string

$$S = \int_{-\infty}^{\infty} dx \int_0^{\hbar\beta} d\tau \left\{ \frac{1}{2} m \left(\frac{\partial u}{\partial \tau} \right)^2 + \frac{\varepsilon_l}{2} \left(\frac{\partial u}{\partial x} \right)^2 + V(u) - \frac{\eta}{2\pi} \frac{\partial u}{\partial \tau} \int_0^{\hbar\beta} d\tau' \frac{\partial u}{\partial \tau'} \ln \left| \sin \frac{\pi}{\hbar\beta} (\tau - \tau') \right| \right\}. \quad (1)$$

Here $\beta = (k_B T)^{-1}$, η is the viscosity coefficient and τ denotes imaginary time. The pinning potential $V(u)$ consists of intrinsic periodic part and the Lorentz potential:

$$V(u) = -\frac{\varphi_0 j_c d}{2\pi} \cos\left(\frac{2\pi u}{d}\right) - \varphi_0 j u, \quad (2)$$

where j_c denotes critical depinning current. For large current, this potential can be expanded around the inflection point to give

$$V(u) = V_0 \left[\left(\frac{u}{w} \right)^2 - \left(\frac{u}{w} \right)^3 \right], \quad (3)$$

where $V_0 = \frac{2}{3} \frac{\varphi_0 j_c^2 \pi^2}{d^2} w^3$ and $w = \frac{3d}{\pi} \left(\frac{j_c - j}{2j_c} \right)^{\frac{1}{2}}$ may be thought as the width of the barrier because $V(0) = V(w) = 0$, and j_c is the critical depinning current. The last term in Eq.(1) is so called Caldeira-Leggett action [12], which describes ohmic damping produced by the coupling to the heat-bath of harmonic oscillators. The line tension ε_l is different for vortices in two different orientations in the **ab** plane. In the semiclassical approximation the decay rate is given by the value of the action on a classical trajectory obtained from the Euler-Lagrange equations of the motion [13].

$$-m \frac{\partial^2 u}{\partial \tau^2} - \varepsilon_l \frac{\partial^2 u}{\partial x^2} + V'(u) + \frac{\eta}{\hbar\beta} \int_0^{\hbar\beta} d\tau' \frac{\partial u}{\partial \tau'} \cot \frac{\pi}{\hbar\beta} (\tau - \tau') = 0 \quad (4)$$

The classical trajectory $u_0(x)$, for static solution of Eq.(4) independent of τ , gives the activation energy U_0 as the static solution of the following equation in the thermal regime $T > T_{cr}$

$$-\varepsilon_l \frac{\partial^2 u_n}{\partial x^2} + V'(u_0) = 0. \quad (5)$$

Below the crossover temperature $T < T_{cr}$ a new kind of trajectory, periodic in imaginary time, develops. Therefore, $u(x, \tau)$ can be expanded in the Fourier series with Matsubara frequencies

$$u(x, \tau) = \sum_{n=0}^{\infty} u_n(x) \cos(\omega_n \tau) \quad ; \quad \omega_n = \frac{2\pi n}{\hbar\beta}. \quad (6)$$

Substituting this expansion into Eq.(3) and linearizing potential around the static solution $u_0(x)$ one obtains

$$-\varepsilon_l \frac{\partial^2 u_n}{\partial x^2} + V''(u_0) u_n = -(\eta \omega_n - m \omega_n^2) u_n. \quad (7)$$

Upon introducing new variables $v_n = \frac{u_n}{w}$ and $\zeta = \frac{x}{d} \left(\frac{\pi^2 w \varphi_0 j_c}{\varepsilon_l} \right)^{\frac{1}{2}}$ the static equation now reads.

$$-\frac{1}{2} \frac{\partial^2 v_0}{\partial \zeta^2} + 2v_0 - 3v_0^2 = 0 \quad (8)$$

Its solution is easily found to be

$$v_0 = \cosh^{-2} \zeta \quad (9)$$

Substitution Eq.(6) into Eq.(4) results in the following equation

$$-\frac{1}{2} \frac{\partial^2 v_n}{\partial \zeta^2} + 2 \left(1 - 3 \cosh^{-2} \zeta\right) v_n = E_n v_n, \quad (10)$$

where

$$E_n = -\frac{j_c w^2}{V_0^2} \left(\eta \omega_n + m \omega_n^2 \right). \quad (11)$$

Eq.9 has three discrete eigenvalues: $-\frac{5}{2}, 0, \frac{3}{2}$ [32]. The negative one determines the crossover temperature

$$k_B T_{cr} = \frac{\hbar \eta}{4\pi m} \left\{ \left[1 + 14.2 \frac{\pi \varphi_0 j_c m}{d \eta^2} \left(1 - \frac{j}{j_c} \right)^{1/2} \right]^{\frac{1}{2}} - 1 \right\} \quad (12)$$

When the static solution of Eq.(5) is used the action S_0 arises only from the elastic and pinning terms in Eq.(1) and the dynamic and inertial terms are cancel out. The action $S_0 = U_0 \hbar \beta$ is then given by:

$$S_0 = 2.77 \left(\frac{\varepsilon_l d^3 \phi_0 j_c}{\pi^3} \right)^{\frac{1}{2}} \left(1 - \frac{j}{j_c} \right)^{5/4} \hbar \beta. \quad (13)$$

Below T_{cr} the predominant mechanism by which the metastable state decays is quantum-mechanical tunneling. The rate at which this takes place is determined by the action for the bounce trajectory. To include contributions of effective mass and viscous damping in the bounce action one would have to go beyond the harmonic approximation in the pinning potential as employed in Eq.(7) and substitute for $u(x, \tau)$ the leading terms in the Matsubara frequencies. Calculations could be performed in the vicinity of crossover temperature in terms of perturbation parameter $(1 - T/T_{cr})$. It can be shown that v_n are of the order $(1 - T/T_{cr})^{\frac{n}{2}}$ and the leading-order contribution in perturbation comes from $v_1 \sim \cosh^{-3} \zeta$. Substituting this term for u in Eq.(1) one can obtain the total action for the bounce trajectory involving contribution of the inertial mass and the viscous damping dynamics.

$$\begin{aligned} S = & 5.26 \hbar \beta \left[\frac{\varepsilon_l d^3 \varphi_0 j_c}{\pi} \right]^{1/2} \left(1 - \frac{j}{j_c} \right)^{5/4} \\ & + \left[\frac{\varepsilon_l d^5}{\pi \varphi_0 j_c} \right]^{1/2} \left(5.71 \frac{m}{\hbar \beta} + 2.85 \frac{\eta}{\pi} \right) \left(1 - \frac{j}{j_c} \right)^{3/4} \end{aligned} \quad (14)$$

2.2. Evaluation of the parameters

The above calculations apply to both kind of vortices which posses different effective masses, line tensions and viscosity coefficients. We shall calculate these parameters as the functions of condensation energy accumulated in the vortex cores. For the stationary flux flow the viscous force $\eta \frac{\partial u}{\partial t}$ is equal to Lorentz force. The electric field generated by

the moving vortex is $E = B \frac{\partial u}{\partial t}$, so we get $E = \frac{\varphi_0 B}{\eta} j = \rho j = \rho_N \frac{B}{H_{c2}} j$ where ρ_N is the normal phase resistivity in the ab plane and H_{c2} is the upper critical field parallel to the layers. Finally,

$$\eta = \frac{\varphi_0 H_{c2}}{\rho_N} = \frac{\varphi_0 \kappa H_c \sqrt{2}}{\rho_N} = \varepsilon_l \frac{4\sqrt{3}\kappa^2}{\pi \rho_N \ln \kappa}, \quad (15)$$

where $H_c = \frac{\varepsilon_l \kappa 2\sqrt{6}}{\pi \varphi_0 \ln \kappa}$ is calculated from the constitutive relation $\varepsilon_l = H_{c1} \varphi_0$.

A moving vortex in the magnetic superconductor can transfer energy to the magnetic moments by emitting spin waves. This energy transfer gives rise to magnetic contribution to the mass and viscosity of the vortex [28]. However the condition for this effect to occur is that the velocity of the vortex should exceed the velocity of the spin wave which is not the case during the flux creep.

The effective mass of the vortex can be deduced from the work of Suhl[31]. He derived the core contribution to the inertial mass $m = \frac{3}{8} m_e \frac{\xi^2 H_c^2 \mu_0}{\epsilon_F}$, where m_e denotes the mass of the electron and ϵ_F the Fermi energy, and the electromagnetic contribution coming from the energy of the electric field induced by the moving flux. Simple estimation shows that this contribution in layered superconductors is 10^{-4} of the core contribution. Therefore,

$$m = \varepsilon_l^2 \frac{9\lambda_{ab}^2 m_e \mu_0}{\varphi_0^2 \pi^2 \epsilon_F (\ln \kappa)^2}. \quad (16)$$

The vortices in two main orientations in the ab plane have different line tension. These ones lying parallel to b direction and those laying in the a direction and created in the magnetic field less then $\frac{1}{2}H_T$ have the line tension ε_b and these ones lying in the a direction but possessing SF domain have the line tension denoted as ε_a

$$\varepsilon_b = \epsilon_0 \ln \frac{\lambda_{ab}}{d}; \quad \varepsilon_a = \frac{\varphi_0 H_T}{2} + \frac{9}{128} \epsilon_0 \ln \frac{\varphi_0 \lambda_c^2}{\pi(\mu_0 H_T + M) d^2 \lambda_{ab}^2}, \quad (17)$$

where $\epsilon_0 = \frac{\varphi_0^2}{4\pi \lambda_{ab} \lambda_c \mu_0}$, M is the magnetic moment of the SF domain inside the vortex and H_T is the thermodynamic critical field of spin-flop domain formation. In order to evaluate the above line tensions it is necessary to calculate first H_T and M . Therefore a following argumentation is proposed. At low fields, in the vicinity of the lower critical field H_{c1} , the intensity of the field in the vortex core is $2H_{c1}$ [30]. When the external field is increased the field intensity in the vortex core increases because of the superposition of the fields of the surrounding vortices. The field intensity in the core must reach H_T in order to originate a transition to the SF phase. Thus, taking into account only z nearest neighbors we can write for the nonunilateral triangular lattice

$$\begin{aligned} \varphi_0 H_T &= 2\varphi_0 H_{c1} + 4z\varepsilon_b \left(\ln \frac{\lambda_{ab}}{d} \right)^{-1} \left[K_0 \left(\frac{c}{\lambda_{ab}} \right) + 2K_0 \left(\frac{c}{2\lambda_{ab}} \sqrt{\frac{3\lambda_c}{\lambda_{ab}}} \right) \right] \\ &= 2\varepsilon_b + o(\varepsilon_b), \end{aligned} \quad (18)$$

where c is the lattice constant. Although there are no precise measurements of spin-flop transition in the antiferromagnetic high temperature superconductors, we assume that

$\mu_0 H_T \approx 40mT$. The typical value of $5.5\mu_B$ per Er ion per unit cell in $ErBa_2Cu_3O_7$ gives $M \approx 0,37$ T and $B_T \approx 0,41$ T. Since $\varphi_0 H_{c1} = \varepsilon_b$ and $d/\xi_c \approx 1$ we obtain:

$$\frac{\varepsilon_a}{\varepsilon_b} = 1 + \frac{36}{128} \frac{\ln \left(\frac{\varphi_0 \lambda_c^2}{\pi(\mu_0 H_T + M) d^2 \lambda_{ab}^2} \right)}{\ln \left(\frac{\lambda_{ab}}{d} \right)} \approx 1.7 \quad (19)$$

It is possible now to relate the viscosity coefficient Eq.(15) and the mass of the vortex Eq.(16) to its line tension Eq.(17).

$$\frac{\eta_a}{\eta_b} = \frac{\varepsilon_a}{\varepsilon_b} \quad (20)$$

and

$$\frac{m_a}{m_b} = \left(\frac{\varepsilon_a}{\varepsilon_b} \right)^2, \quad (21)$$

With the use of a simple consideration we can estimate the change of j_c due to the creation of spin-flop domain along the vortex

$$j_c \varphi_0 d \sim \frac{1}{2} \int dy dz C(y, z) \left(\frac{\partial u_z}{\partial z} \right)^2,$$

where the Fourier transform of the compression modulus is given by [33]

$$C(k_y, k_z) = \frac{B^2}{\mu_0(1 + \lambda_{ab}^2 k_z^2 + \lambda_c^2 k_y^2)}$$

By taking $dy dz \sim \varphi_0/B$; $u_z \sim d$; $\frac{\partial}{\partial z} \sim k_z \sim k_y(\lambda_c/\lambda_{ab})$ the integral can be estimated as follows

$$j_c = \frac{Bd}{4\lambda_{ab}^2}. \quad (22)$$

In the a direction, however, we have an additional contribution from the magnetic domain $dy dz \sim 5\varphi_0/8\pi B_T$, so we get

$$j_{ca} = j_c + \frac{5dB_T}{128\lambda_{ab}^2} \quad (23)$$

and finally we get $j_{ca} \approx 3j_{cb}$.

3. Motion of the flux in the quantum regime

Consider now a hollow cylindrical sample of the radius a and the wall thickness $g \ll a$ (this is so called thin wall approximation). The sample is placed in a magnetic field directed parallel to the axis of the cylinder and along the superconducting layers. A trapped flux in the system is $\Phi \cong (B_{in} - B_{ex})$, where B_{in} denotes the field inside the hole of the sample and B_{ex} outside the sample respectively. The motion of the flux is triggered off by an activation process in which segments of the flux line tunnel through an intrinsic pinning potential to the neighboring interlayer spacing. By applying Faraday's

law we can easily calculate electric field in the sample due to the change of the trapped flux

$$E = -\frac{1}{2}\mu_0 a g j_c \frac{d(j/j_c)}{dt}$$

which is equal to the mean electric field associated with the motion of vortices $E = \varphi_0 W L d$, where L is the length of the sample in the direction of the applied field, d interlayer separation and W is the activation probability per unit volume and unit time. The probability for quantum tunneling is given by $W \approx \exp(-S/\hbar)$. As was shown by Tekiel [34] the activation probability of macroscopic quantum excitations is proportional to j^3 . Thus we can assume that $W = \alpha_0 j^3 \exp(-S/\hbar)$. Combining all above together we get

$$\Omega t = - \int_{x(t)}^{x(0)} \frac{\exp(-S(x)/\hbar)}{x^3} dx \quad (24)$$

where $x = j/j_c$, $\Omega = \varphi_0 \alpha_0 j_c^2 / (\mu_0 \gamma)$, $\gamma = \frac{a g}{L d}$ is a geometrical factor determined by the shape of a sample. Equation (21) was solved numerically for the input parameters shown in Tab.(1), and the results are shown in the Fig.(5). As we can see the creep is

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|--------------------------|--------------------------------------|
| $d = 10^{-9} m$ | $\eta_b = 10^{-8} kg/(s \cdot m)$ |
| $j_{cb} = 10^{12} A/m^2$ | $\lambda_{ab} = 25 \times 10^{-9} m$ |
| $m_b = 10^{-22} kg/m$ | $\lambda_c = 125 \times 10^{-9} m$ |
| $\epsilon_F = 0.1 eV$ | $\mu_0 H_{c2}^{ab} = 150 T$ |

Table 1. Input parameters

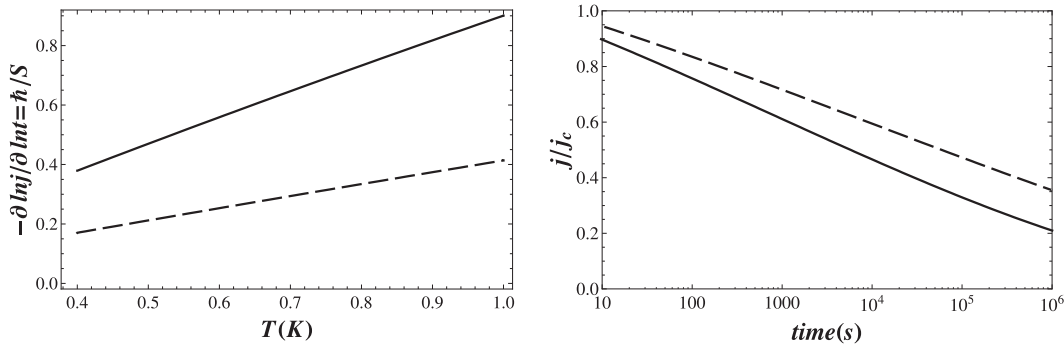


Figure 2. Left panel: Normalized creep rate as a function of the temperature. Right panel: Quantum creep rate (the rate of the decay of the trapped flux $\Phi/\Phi_0 \sim j/j_c$) as a function of the logarithm of time. Dashed line depicts vortices possessing magnetic cores.

slower when vortices possessing magnetic domain occurred in the system.

4. Discussion of the results

The main purpose of the present study is to derive expressions for the bounce contribution to the action arising from the inertial mass term, elastic and pinning

terms, and the damping term. As was explained previously, in antiferromagnetic superconductors it is possible that two distinct types of vortices occur in the sample. One that is quite similar to the vortices of nonmagnetic superconductor and the other possessing a ferromagnetic-like domain around, and inside, the core. What distinguishes both types of vortices in the present calculations is their tension Eq.(18) and critical depinning current Eq.(22). We consider the limit of a large current which is only slightly less than critical depinning current. In this approximation we numerically calculated the crossover temperature from Eq.(12) and plotted it as a function of (FC), fractional current j/j_c

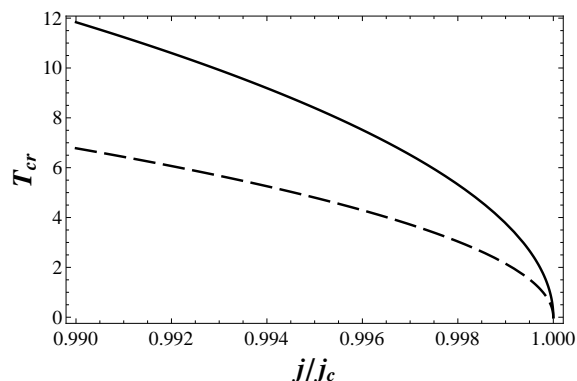


Figure 3. Crossover temperature as a function of fractional current j/j_c . Dashed line depicts vortices possessing magnetic cores.

As can be seen the difference of crossover temperatures for both types of vortices depends on the critical current, and for the value of the FC equal to 0.99 this difference reaches almost 5 K, and then vanishes when FC reaches 1. Fig.(4) shows the activation energy U_0 calculated from Eq.(13) as a function of FC. The activation energy differs about 1.5 meV for 0.99 of FC and goes to 0 when FC reaches 1.

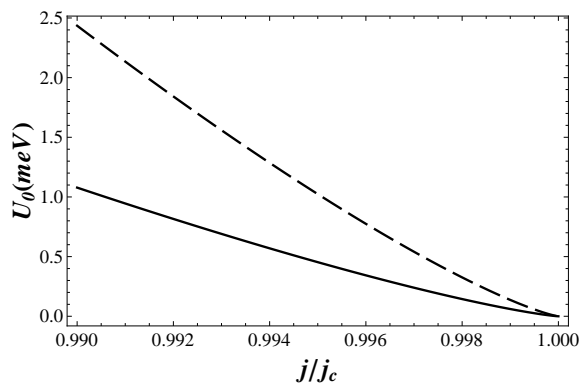


Figure 4. Activation energy U_0 (in meV) calculated as a function of fractional current j/j_c . Dashed line depicts vortices possessing magnetic cores.

Another important effect is that the action and hence activation energy is rendered temperature dependent so that the quantum tunneling rate becomes temperature

dependent below the crossover temperature. Moreover, including the damping and inertial mass contributions reduces the quantum tunneling rate. From the above considerations one can draw yet another conclusion. At constant temperature the flux creep regime can be altered. If the direction or intensity of the applied magnetic field is changed, it is possible to leap from quantum to thermal creep, and vice versa. Let

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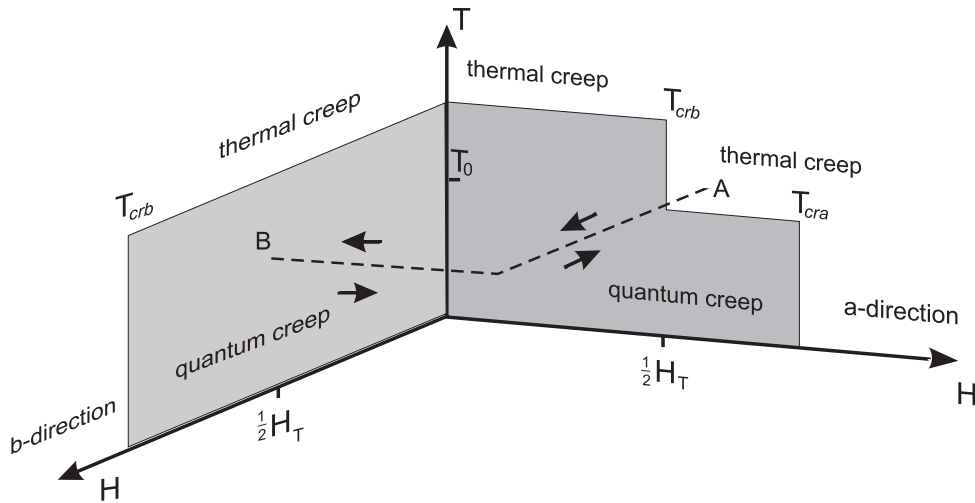


Figure 5. Schematic diagram of the leap from thermal to quantum creep, and vice versa.

us see how it can be performed. Let us increase magnetic field intensity, aligned in the **a** direction, to the point marked A. The system is in the thermal creep regime if the temperature of the sample is fixed somewhere in the range $T_{crb} > T_0 > T_{cra}$, as shown in Fig.(5). Then we change the direction of external field from **a** to **b** axis. The system now leaps over to the point B and finds itself in the quantum creep regime. Doing the same operation in the reverse order one enforces the system to crossover from quantum to thermal creep. Yet another scenario is possible. When magnetic field is applied along **a** axis and the temperature is kept constant in the interval $T_{crb} > T_0 > T_{cra}$, the increase of magnetic field intensity above $\frac{1}{2}H_T$ enforces new vortices to occur, and the leap of the system from a quantum to thermal creep regime. This scenario could easily be named a "creep valve". It enables to make the creep rate slower or faster simply by manipulation of the magnetic field intensity around its value $\frac{1}{2}H_T$.

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